

Recent Progress in Nonperturbative QCD Theory and Insight on Cosmological Phase Transition

H. Suganuma, H. Ichie, H. Toki and H. Monden*

*Research Center for Nuclear Physics (RCNP), Osaka University
Mihogaoka 10-1, Ibaraki 567, Osaka, Japan*

** Department of Physics, Tokyo Metropolitan University
Minami-Osawa 1-1, Hachioji 192, Tokyo, Japan*

The QCD phase transition at finite temperature is studied with the dual Ginzburg-Landau theory, which is the QCD effective theory based on the dual Higgs mechanism by QCD-monopole condensation. At high temperature, the confinement force is largely reduced by thermal effects, which leads to the swelling of hadrons. Simple formulae for the surface tension and the thickness of the phase boundary are derived from the shape of the effective potential at the critical temperature. We investigate also the process of the hadron-bubble formation in the early Universe.

1. Introduction

Quantum chromodynamics (QCD) is the fundamental theory of the strong interaction [1-3]. In spite of the simple form of the QCD lagrangian,

$$\mathcal{L}_{\text{QCD}} = -\frac{1}{2}\text{tr}G_{\mu\nu}G^{\mu\nu} + \bar{q}(\not{D} - m_q)q, \quad (1.1)$$

it miraculously provides quite various phenomena like color confinement, dynamical chiral-symmetry breaking, non-trivial topologies, quantum anomalies and so on, as shown in Fig.1. Then, it would be interesting to compare QCD with the history of the Universe, because a quite simple ‘big bang’ also created various things including galaxies, stars, lives and thinking reeds [4]. Therefore, QCD can be regarded as an interesting miniature of the history of the Universe. This is the most attractive point of the QCD physics.

Since it is quite difficult to understand the various QCD phenomena and their underlying mechanism at the same time, many methods and models have been proposed to understand each phenomenon. We show in Fig.2 a brief sketch on the history of QCD and typical QCD effective models [1] as listed in Table 1. The bag model and the Nambu-Jona-Lasinio (NJL) model are simple relativistic quark theories. The nonlinear σ model and the Skyrme-Witten soliton model are described by the Nambu-Goldstone pion. In particular, the Skyrme-Witten model is stimulative because a fermionic baryon can be described as a soliton made by bosonic pions in this framework. The origin of this magic is found in the non-trivial topology of the flavor dynamics in QCD. The lattice QCD simulation [5] using a supercomputer is a hopeful and promising method based on QCD directly, and its importance becomes larger and larger according to the great progress of the computational power. In particular, recent lattice QCD studies [6] shed a light

on the confinement mechanism, which is one of the most difficult problems in the particle physics. In these years, the origin of color confinement can be recognized as the dual Higgs mechanism by monopole condensation, and the nonperturbative QCD vacuum is regarded as the dual superconductor. The dual Ginzburg-Landau (DGL) theory [7-14] was formulated with this picture, and can reproduce confinement properties like the string tension and the hadron flux-tube formation.

Now, you may find a current of the QCD physics. In '80s, chiral symmetry breaking was the central issue. The chiral bag model, the NJL model and the σ model were formulated with referring chiral symmetry. In '90s, on the other hand, the confinement physics is providing an important current of the hadron physics. The key word for the understanding of confinement is the “duality”, which is recently paid attention by many theoretical particle physicists after Seiberg-Witten’s discovery on the essential role of monopole condensation for the confinement in a supersymmetric version of QCD [15].

2. Color Confinement and Dual Higgs Mechanism

We briefly show the modern current of the confinement physics. About 20 years ago, Nambu, 't Hooft and Mandelstam proposed an interesting picture for color confinement based on the analogy between the superconductor and the QCD vacuum [16]. In the superconductor, the magnetic field is excluded due to the Meissner effect, which is caused by Cooper-pair condensation. As the result, the magnetic flux is squeezed like the Abrikosov vortex. On the other hand, the color-electric flux is excluded in the QCD vacuum, and therefore the squeezed color-flux tube is formed between color sources. Thus, these two systems are quite similar and can be regarded as the dual version each other. This idea is based on the “duality” of the gauge theories, which was firstly pointed out by Dirac more than 50 years ago [17].

With referring Table 2 and Fig.3, we compare the ordinary electromagnetic system, the superconductor and the nonperturbative QCD vacuum regarded as the dual superconductor. In the ordinary electromagnetism in the Coulomb phase, both electric flux and magnetic flux are conserved, respectively. The electric-flux conservation is guaranteed by the ordinary gauge symmetry. On the other hand, the magnetic-flux conservation is originated from the dual gauge symmetry [8-14], which is the generalized version of the Bianchi identity. As for the inter-charge potential in the Coulomb phase, both electric and magnetic potentials are Coulomb-type.

The superconductor in the Higgs phase is characterized by electric-charge condensation, which leads to the Higgs mechanism or spontaneous breaking of the ordinary gauge symmetry, and therefore the electric flux is no more conserved. In such a system obeying the London equation, the electric inter-charge potential becomes short-range Yukawa-type similarly in the electro-weak unified theory. On the other hand, the dual gauge symmetry is not broken, so that the magnetic flux is conserved, but is squeezed like a one-dimensional flux tube due to the Meissner

effect. As the result, the magnetic inter-charge potential becomes linearly rising like a condenser.

The nonperturbative QCD vacuum regarded as the dual Higgs phase is characterized by color-magnetic monopole condensation, and resembles the dual version of the superconductor, where the “dual version” means the interchange between the electric and magnetic sectors. Monopole condensation leads to the spontaneous breaking of the dual gauge symmetry, so that color-magnetic flux is not conserved, and the magnetic inter-charge potential becomes short-range Yukawa-type. Note that the ordinary gauge symmetry is not broken by such monopole condensation. Therefore, color-electric flux is conserved, but is squeezed like a one-dimensional flux-tube or a string as a result of the dual Meissner effect. Thus, the hadron flux-tube is formed in the monopole-condensed QCD vacuum, and the electric inter-charge potential becomes linearly rising, which confines the color-electric charges as quarks [8-11].

As a remarkable fact in the duality physics, these are two “see-saw relations” between in the electric and magnetic sectors.

- (1) There appears the Dirac condition $eg = 2\pi$ [17] in QED or $eg = 4\pi$ [8] in QCD. Here, unit electric charge e is the gauge coupling constant, and unit magnetic charge g is the dual gauge coupling constant. Therefore, a strong-coupling system in one sector corresponds to a weak-coupling system in the other sector.
- (2) As shown in Fig.3, the long-range confinement system in one sector corresponds to a short-range (Yukawa-type) interaction system in the other sector.

Let us consider usefulness of the latter “see-saw relation”. One faces highly non-local properties among the color-electric charges in the QCD vacuum because of the long-range linear confinement potential. Then, the direct formulation among the electric-charged variables would be difficult due to the non-locality, which seems to be a destiny in the long-distance confinement physics. However, one finds a short-range Yukawa potential in the magnetic sector, so that the electric-confinement system can be approximated by a local formulation among magnetic-charged variables. Thus, the confinement system, which seems highly non-local, can be described by a short-range interaction theory using the dual variables, which is the most attractive point in the dual Higgs theory.

Color-magnetic monopole condensation is necessary for color confinement in the dual Higgs theory. As for the appearance of color-magnetic monopoles in QCD, 't Hooft proposed an interesting idea of the abelian gauge fixing [18,19], which is defined by the diagonalization of a gauge-dependent variable. In this gauge, QCD is reduced into an abelian gauge theory with the color-magnetic monopole [8,18], which will be called as QCD-monopoles hereafter. Similar to the 't Hooft-Polyakov monopole [2] in the Grand Unified theory (GUT), the QCD-monopole appears from a hedgehog configuration corresponding to the non-trivial homotopy group $\pi_2(\text{SU}(N_c)/\text{U}(1)^{N_c-1}) = \mathbb{Z}_{\infty}^{N_c-1}$ on the nonabelian manifold.

Many recent studies based on the lattice gauge theory have supported QCD-monopole condensation and abelian dominance [6,13,20-22], which means that only abelian variables are relevant for the nonperturbative QCD, in the 't Hooft abelian

gauge. Hence, the dual-superconductor scenario seems workable for color confinement in the QCD vacuum, and the nonperturbative QCD would be described by the dual Ginzburg-Landau (DGL) theory, which is the QCD effective theory based on the dual Higgs mechanism.

3. Dual Ginzburg-Landau Theory

The dual Ginzburg-Landau (DGL) lagrangian [13,14] for the pure-gauge system is described with the dual gauge field B_μ and the QCD-monopole field χ ,

$$\mathcal{L}_{\text{DGL}} = \text{tr} \left\{ -\frac{1}{2}(\partial_\mu B_\nu - \partial_\nu B_\mu)^2 + [\hat{D}_\mu, \chi]^\dagger [\hat{D}^\mu, \chi] - \lambda(\chi^\dagger \chi - v^2)^2 \right\}, \quad (3.1)$$

where $\hat{D}_\mu \equiv \partial_\mu + igB_\mu$ is the dual covariant derivative including the dual gauge coupling constant $g = 4\pi/e$.

The dual gauge field $B_\mu \equiv \vec{B}_\mu \cdot \vec{H} = B_\mu^3 T^3 + B_\mu^8 T^8$ is defined on the dual gauge manifold $U(1)_m^3 \times U(1)_m^8$ [13,14], which is the dual space of the maximal torus subgroup $U(1)_e^3 \times U(1)_e^8$ embedded in the original gauge group $SU(3)_c$. The abelian field strength tensor is written as $F_{\mu\nu} = *(\partial \wedge B)_{\mu\nu}$ so that the role of the electric and magnetic fields are interchanged in comparison with the ordinary A_μ description.

The QCD-monopole field χ is defined as $\chi \equiv \sqrt{2} \sum_{\alpha=1}^3 \chi_\alpha E_\alpha$ [13,14] where $E_1 \equiv \frac{1}{\sqrt{2}}(T_6 + iT_7)$, $E_2 \equiv \frac{1}{\sqrt{2}}(T_4 - iT_5)$ and $E_3 \equiv \frac{1}{\sqrt{2}}(T_1 + iT_2)$. Here, χ_α has the magnetic charge $g\vec{\alpha}$ proportional to the root vector $\vec{\alpha}$. In the QCD-monopole condensed vacuum with $|\chi_\alpha| = v$, the dual gauge symmetry $U(1)_m^3 \times U(1)_m^8$ is spontaneously broken instead of the gauge symmetry $U(1)_e^3 \times U(1)_e^8$. Through the dual Higgs mechanism, the dual gauge field B_μ acquires its mass $m_B = \sqrt{3}gv$, whose inverse provides the radius of the hadron flux tube [8], and the dual Meissner effect causes the color-electric field excluded from the QCD vacuum, which leads to color confinement. The QCD-monopole fluctuations $\tilde{\chi}_\alpha \equiv \chi_\alpha - v$ ($\alpha=1,2,3$) also acquire their mass $m_\chi = 2\sqrt{\lambda}v$ in the QCD-monopole condensed vacuum. As a relevant prediction, only one QCD-monopole fluctuation $\tilde{\chi} \equiv \sum_{\alpha=1}^3 \tilde{\chi}_\alpha$ appears as a color-singlet scalar glueball in the confinement phase, although the dual gauge field B_μ and the other two combinations of the QCD-monopole fluctuation are not color-singlet and cannot be observed [13,21,22].

The DGL theory reproduces confinement properties like the inter-quark potential and the hadron flux-tube formation. We studied effects of the flux-tube breaking by the light quark-pair creation in the DGL theory, and derived the infrared screened inter-quark potential [8-10], which is observed in the lattice QCD with dynamical quarks [5]. We studied also the dynamical chiral-symmetry breaking (D χ SB), which is also an important feature in the nonperturbative QCD, by solving the Schwinger-Dyson equation for the dynamical quark [8-10]. The quark-mass generation is brought by QCD-monopole condensation, which suggests the

close relation between $D\chi SB$ and color confinement. Thus, the DGL theory provides not only the confinement properties but also $D\chi SB$ and its related quantities like the constituent quark mass, the chiral condensate and the pion decay constant [8-10,23,24].

4. QCD Phase Transition in the DGL Theory

In this chapter, we study the QCD phase transition [10-12,14,24] in the DGL theory. Although quarks and gluons are confined inside hadrons in the nonperturbative QCD vacuum, these colored particles are liberated and the system becomes the quark-gluon-plasma (QGP) phase [1] at high temperature ($T > T_c \sim 200\text{MeV}$), which is called as the QCD phase transition. The experimental creation of the QGP, one of the most central subjects in the RHIC project, is expected to be realized in ultra-relativistic heavy-ion collisions [1,25]. On the other hand, the QCD phase transition occurred in the early Universe, and its process strongly influenced the afterward nucleosynthesis [26].

In the DGL theory, the QCD phase transition is characterized by QCD-monopole condensate $\bar{\chi} \equiv |\chi_\alpha|$, which is an order parameter on the confinement strength. The QCD-monopole condensed vacuum ($\chi \neq 0$), where the dual gauge symmetry is spontaneously broken ($m_B \neq 0$), is the confinement phase with a non-vanishing string tension ($k \neq 0$). Without QCD-monopole condensation ($\chi = 0$), the dual gauge symmetry is manifest ($m_B = 0$), and the system corresponds to the deconfinement phase, where the string tension disappears ($k = 0$).

To study the QCD phase transition, the effective potential $V_{\text{eff}}(\bar{\chi}; T)$ at finite temperature, which physically means the thermodynamical potential, is formulated as the function of QCD-monopole condensate $\bar{\chi} \equiv |\chi_\alpha|$ [10-12,14],

$$V_{\text{eff}}(\bar{\chi}; T) = 3\lambda(\bar{\chi}^2 - v^2)^2 + 3T \int \frac{d^3k}{(2\pi)^3} [2\ln(1 - e^{-\sqrt{k^2 + m_B^2}/T}) + \ln(1 - e^{-\sqrt{k^2 + m_\chi^2}/T})] \quad (4.1)$$

with $m_B = \sqrt{3}g\bar{\chi}$ and $m_\chi = 2\sqrt{\lambda}\bar{\chi}$. Here, we have used the quadratic source term to avoid the imaginary scalar-mass problem [10-12,14].

The QCD-monopole condensate $\bar{\chi}_{\text{phys}}(T)$ at finite temperature is obtained from the local minimum of the effective potential $V_{\text{eff}}(\bar{\chi}; T)$. As temperature increases, $\bar{\chi}_{\text{phys}}(T)$ decreases and disappears at a critical temperature T_c . Above T_c , the deconfinement phase is realized as the Coulomb phase with $\bar{\chi}_{\text{phys}}(T) = 0$, where the dual gauge symmetry is restored. With the parameters in Ref.[8], we find a weak first order phase transition at $T_c = 0.2\text{GeV}$, and the mixed phase or the two-phase coexistence is allowed only in $T_{\text{low}} < T < T_{\text{up}}$ with $T_{\text{low}} \simeq 0.193\text{GeV}$ and $T_{\text{up}} \simeq 0.201\text{GeV}$ [10,12,14].

Now, let us consider the confinement and hadron properties. Fig.4 shows the string tension $k(T)$ at finite temperature. As the temperature increases, $k(T)$ becomes smaller and drops rapidly near T_c [10-12,14,27]. Therefore, the slope of the inter-quark potential is reduced and the inter-quark distance inside hadrons

increases at high temperature. In addition, the color-electric field spreads according to the decrease of m_B . Thus, the reduction of the confinement force leads to the swelling of hadrons at high temperature [14].

We predict also a large mass reduction of the scalar glueball originated from QCD-monopoles near T_c [10-12,14]. We guess that the violent excitation of these scalar glueballs with a reduced mass would promote the QCD phase transition.

Finally, we consider the relation between the surface tension σ and the effective potential $V_{\text{eff}}(\bar{\chi}; T_c)$ [10,11]. The surface tension σ characterizes the strength of the first order in the phase transition, and is very important for the shape of the boundary surface in the mixed phase, where the two phases correspond to the two minima, $\bar{\chi} = 0$ and $\bar{\chi} = \bar{\chi}_c$, in $V_{\text{eff}}(\bar{\chi}; T_c)$. Using the sine-Gordon kink ansatz [10,11] for the boundary profile, $\bar{\chi}(z) = \bar{\chi}_c \tan^{-1} e^{z/\delta}$, we derive simple formulae for the surface tension $\sigma \simeq \frac{4\sqrt{3}}{\pi} \sqrt{h} \bar{\chi}_c$, and the phase-boundary thickness $2\delta \simeq \frac{2\sqrt{3}}{\pi} \bar{\chi}_c / \sqrt{h}$, where h is the “barrier height” between the two minima in $V_{\text{eff}}(\bar{\chi}; T_c)$. We find $2\delta \simeq 3.4\text{fm}$ and $\sigma \simeq (112\text{MeV})^3$ [10], which seems consistent with the lattice QCD data [28].

5. Hadron Bubble Formation in the Early Universe

In this chapter, we study the hadron bubble formation [29] in the early Universe using the DGL theory [14]. As Witten pointed out [30], if the QCD phase transition is of the first order, the hadron and QGP phases should coexist in the early Universe. During the mixed-phase period, there appears the inhomogeneity on the baryon density distribution [26,31], which can strongly affects the primordial nucleosynthesis and the successive history of the Universe.

Let us consider how hadron bubbles appear in the QGP phase near the critical temperature T_c in the DGL theory. The hadron bubbles are created in the super-cooling QGP phase with $T_{\text{low}} < T < T_c$. We use the sine-Gordon kink ansatz [14] for the profile of the QCD-monopole condensate in the hadron bubble as a function of the radial coordinate r , $\bar{\chi}(r; R) = \bar{\chi}_{\text{phys}}(T) \tan^{-1} e^{(R-r)/\delta} / \tan^{-1} e^{R/\delta}$, where R and 2δ correspond to the hadron-bubble radius and the phase-boundary thickness, respectively. The QCD-monopole condensate $\bar{\chi}(r; R)$ is finite only inside the bubble, $r \lesssim R$. The energy density $\mathcal{E}(r; R)$ of the hadron bubble is shown in Fig.5. It is negative inside and positive near the boundary surface. The total energy is roughly estimated as the sum of the positive surface term and the negative volume term.

The total energy of the hadron bubble with radius R can be estimated using the effective potential $V_{\text{eff}}(\bar{\chi}; T)$, $E(R; T) = 4\pi \int_0^\infty dr r^2 \{3(\frac{d\bar{\chi}(r; R)}{dr})^2 + V_{\text{eff}}(\bar{\chi}; T)\}$, where the thickness δ is determined by the energy minimum condition. The hadron-bubble energy $E(R; T)$ as shown in Fig.6 takes a maximal value at a critical radius R_c , which leads the collapse of the hadron bubbles with smaller radius than R_c . Only large hadron bubbles with $R > R_c$ grow up with radiating the shock wave [14,29].

On the other hand, the creation of large bubbles is suppressed because of the small creation probability [14]. In the hadron-bubble formation, there is a penetration over a large barrier height $h(T)$ in the effective potential per unit volume, and therefore the creation of large bubbles needs a large energy fluctuation. Such a process is strongly suppressed because of the thermodynamical factor $P(T) \equiv \exp(-\frac{4\pi}{3}R_c(T)^3h(T)/T)$, which is proportional to the hadron-bubble formation rate. Thus, the only small bubbles are created practically, although its radius should be larger than R_c [14].

Using $V_{\text{eff}}(\bar{\chi}; T)$ in the DGL theory, we can estimate the critical radius $R_c(T)$ and the hadron-bubble formation factor $P(T)$ at finite temperature as shown in Fig.7 and Fig.8, respectively. As the temperature decreases, the created hadron bubbles becomes smaller, while the bubble formation rate becomes larger [14].

From these results, we can imagine how the QCD phase transition happens in the big bang scenario [14] as shown in Fig.9.

- (a) Slightly below T_c , only large hadron bubbles appear, but the creation rate is quite small.
- (b) As temperature is lowered by the expansion of the Universe, smaller bubbles are created with much formation rate. During this process, the created hadron bubbles expand with radiating shock wave, which reheats the QGP phase around them [29].
- (c) Near T_{low} , many small hadron bubbles are violently created in the unaffected region free from the shock wave.
- (d) The QGP phase pressured by the hadron phase is isolated as high-density QGP bubbles [29,30], which provide the baryon density fluctuation [26].

Thus, the numerical simulation using the DGL theory would tell how the hadron bubbles appear and evolve quantitatively in the early Universe.

We would like to thank Prof. Kajino for fruitful discussions on the hadron bubble formation.

REFERENCES

1. W. Greiner and A. Schäfer, “Quantum Chromodynamics”, (Springer, 1994).
2. T. P. Cheng and L. F. Li, “Gauge Theory of Elementary Particle Physics” (Clarendon press, Oxford, 1984).
3. K. Huang, “Quarks, Leptons and Gauge Fields”, (World Scientific, 1982).
4. B. Pascal, “Pensées” (1670).
5. H. J. Rothe, “Lattice Gauge Theories”, (World Scientific, 1992).
6. A. S. Kronfeld, G. Schierholz, U. -J. Wiese, Nucl. Phys. **B293** (1987) 461.
S. Hioki, S. Kitahara, S. Kiura, Y. Matsubara, O. Miyamura, S. Ohno and T. Suzuki, Phys. Lett. **B272** (1991) 326.
S. Kitahara, Y. Matsubara and T. Suzuki, Prog. Theor. Phys. **93** (1995) 1.

7. T. Suzuki, Prog. Theor. Phys. **80** (1988) 929 ; **81** (1989) 752.
S. Maedan and T. Suzuki, Prog. Theor. Phys. **81** (1989) 229.
8. H. Suganuma, S. Sasaki and H. Toki, Nucl. Phys. **B435** (1995) 207.
9. H. Suganuma, S. Sasaki and H. Toki, Proc. of Int. Conf. on “Quark Confinement and Hadron Spectrum”, Como Italy, (World Scientific, 1995) 238.
10. H. Suganuma, S. Sasaki, H. Toki and H. Ichie, Prog. Theor. Phys. (Suppl.) **120** (1995) 57.
11. H. Suganuma, H. Ichie, S. Sasaki and H. Toki, “Color Confinement and Hadrons”, (World Scientific, 1995) 65.
12. H. Ichie, H. Suganuma and H. Toki, Phys. Rev. **D52** (1995) 2944.
13. H. Suganuma, S. Umisedo, S. Sasaki, H. Toki and O. Miyamura, Proc. of “Quarks, Hadrons and Nuclei”, Adelaide Australia, Nov. 1995, to appear in Aust. J. Phys.
14. H. Ichie, H. Monden, H. Suganuma and H. Toki, Proc. of “Nuclear Reaction Dynamics of Nucleon-Hadron Many Body Dynamics”, Osaka, Dec. 1995 (World Scientific) in press.
15. N. Seiberg and E. Witten, Nucl. Phys. **B426** (1994) 19; **B431** (1994) 484.
G. P. Collins, Physics Today, March (1995) 17.
16. Y. Nambu, Phys. Rev. **D10** (1974) 4262.
G. 't Hooft, “High Energy Physics” (Editorice Compositori, Bologna 1975).
S. Mandelstam, Phys. Rep. **C23** (1976) 245.
17. P. A. M. Dirac, Proc. Roy. Soc. **A133** (1931) 60.
18. G. 't Hooft, Nucl. Phys. **B190** (1981) 455.
19. Z. F. Ezawa and A. Iwazaki, Phys. Rev. **D25** (1982) 2681; **D26** (1982) 631.
20. O. Miyamura, Nucl. Phys. **B**(Proc. Suppl.)**42** (1995) 538.
O. Miyamura and S. Origuchi, *Color Confinement and Hadrons*, (World Scientific, 1995) 235.
21. H. Suganuma, A. Tanaka, S. Sasaki and O. Miyamura, Proc. of “Lattice Field Theories '95”, Nucl. Phys. **B** (Proc. Suppl.) **47** (1996) 302.
22. H. Suganuma, K. Itakura, H. Toki and O. Miyamura, Proc. of “Non-perturbative Approaches to QCD”, Trento Italy, July 1995 (PNPI, 1995) 224.
23. S. Sasaki, H. Suganuma and H. Toki, Prog. Theor. Phys. **94** (1995) 373.
24. S. Sasaki, H. Suganuma and H. Toki, Proc. of Int. Conf. on “Baryons '95”, Santa Fe, Oct. 1995 in press.
25. H. Ichie, H. Suganuma and H. Toki, preprint, Phys. Rev. **D** in press.
26. T. Kajino and R. N. Boyd, Ap. J. **359** (1990) 267.
T. Kajino, G. J. Mathew and G. M. Fuller, Ap. J. **364** (1990) 7.
27. M. Gao, Nucl. Phys. **B9** (Proc. Suppl.) (1989) 368.
28. Y. Iwasaki, K. Kanaya, L. Karkkainen, Phys. Rev. **D49** (1994) 3540.

- 29. G. M. Fuller, G. J. Mathews, C. R. Alcock, Phys. Rev. **D 37** (1988) 1380.
- 30. E. Witten, Phys. Rev. **D30** (1984) 272.
- 31. T. Kajino, Phys. Rev. Lett. **66** (1991) 125.
T. Kajino, M. Orito, Y. Yamamoto and H. Suganuma, “Color Confinement and Hadrons”, (World Scientific, 1995) 263.

Table 1: Features of the lattice QCD and the QCD effective models.

Table 2: The comparison among the ordinary electromagnetic system, the superconductor and the nonperturbative QCD vacuum regarded as the QCD-monopole condensed system.

Fig.1: A brief sketch of the QCD physics.

Fig.2: A brief history of QCD and typical QCD effective models.

Fig.3 : The electric and magnetic inter-charge potentials in the Coulomb, Higgs and dual Higgs phases. The “see-saw relation” is found between in the electric and magnetic sectors. The long-range confinement system in one sector corresponds to a short-range (Yukawa-type) interaction system in the other sector.

Fig.4: The string tension $k(T)$ at finite temperature T . The black dots denote lattice QCD data [27] in the pure gauge.

Fig.5: The energy density $\mathcal{E}(r; R)$ of the hadron bubble. It is negative inside and positive near the boundary surface.

Fig.6: The energy of the hadron bubble $E(R; T)$ as function of the hadron bubble radius R .

Fig.7: The critical radius $R_c(T)$, which provides maximum of $E(R; T)$, at finite temperature T .

Fig.8: The hadron-bubble formation factor $P(T) \equiv \exp(-\frac{4\pi}{3}R_c^3(T)h(T)/T)$ as a function of temperature T .

Fig.9: The appearance and evolution of hadron bubbles in the early Universe. The shaded and white regions denote the hadron and QGP phases, respectively. (a) Slightly below T_c , only large hadron bubbles appear with very small creation rate. (b) Hadron bubbles expand with radiating shock wave, which reheats the QGP phase around them. The affected region is expressed by the arrow. (c) Near T_{low} , many small hadron bubbles are violently created in the unaffected region. (d) The QGP phase pressured by the hadron phase is isolated as high-density QGP bubbles.